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Nucleon and $\gamma N \to \Delta$ lattice form factors in a constituent quark model

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A covariant quark model, based both on the spectator formalism and on vector meson dominance, and previously calibrated by the physical data, is here extended to the unphysical region of the lattice data by means of one single extra adjustable parameter – the constituent quark mass in the chiral limit. We calculated the Nucleon (N) and the $\gamma N \to \Delta$ form factors in the universe of values for that parameter described by quenched lattice QCD. A qualitative description of the Nucleon and $\gamma N \to \Delta$ form factors lattice data is achieved for light pions.

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I. INTRODUCTION

In principle QCD is the fundamental theory that describes hadronic systems, but its non-perturbative character makes a direct application to hadrons difficult, with the exception of the high energy and momentum transfer regime. Fortunately, at low Q^2 effective field theories based on effective hadronic degrees of freedom have been applied with success. Examples are chiral perturbation theory [1], small scale expansion [2], soliton models and constituent quark models [3, 4, 5]. Constituent quark models provided consistent descriptions encompassing the low-energy baryon spectrum, elastic and inelastic form factors, charge and magnetic radii, magnetic moments, axial and pseudoscalar form factors [6, 7, 8, 9, 10, 11, 12].

Although all those different frameworks may include the essential features of QCD, which are confinement and chiral symmetry, they are only partial simulations of the true underlying theory.

Recently, significant progress has occurred in calculations of QCD in the lattice, which evaluate the important QCD non-perturbative contributions at low Q^2 directly from the underlying theory. Up to now the applications of lattice QCD are restricted still to large pion masses ($m_\pi > 350$ MeV) corresponding to heavy quarks, and lattice spacings that are one order of magnitude smaller than the size of the nucleon ($a \sim 0.05$ fm). To extract information on the real world the results must be extrapolated both to the continuum limit ($a \to 0$) and to the physical pion mass regime [1, 13]. However, except for the limit to the physical region, there is no simple way of interpreting the lattice QCD data.

In this work we invert this procedure. We take here the challenge of describing the lattice data using a quark model. We start with the constituent quark model obtained from the covariant spectator theory. Our model was presented in Refs. [4, 14], where it was adjusted to the experimental data for the four nucleon elastic form factors and the dominant form factor of the $\gamma N \to \Delta$ transition. In this work we extend this model to the unphysical region of the lattice data. Although the constituent quark models were originally thought for and applied to the physical limit, a constituent quark model that is simultaneously consistent with the experimental data and lattice QCD data is valuable, because it includes indirectly the constraints of QCD, and therefore satisfies properties of the underlying theory. A similar procedure was considered in Refs. [15, 16] using a larger number of parameters.

Our model does not include explicit pion cloud effects. However, as the electromagnetic interaction with the quarks is parametrized according to vector meson dominance (VMD), part of the pion cloud effects are indirectly taken into account. Therefore the model can be applied in the regions where the valence quark degrees of freedom are dominant, and it gives a good description of the nucleon and Δ systems [14, 17].

For the nucleon, pion cloud effects are fundamental in the time-like region [18], although they are not so important in the region of small positive Q^2 values [14], which is involved in the study of the electromagnetic transition. For the $\gamma N \to \Delta$ transition, as it happens also to other constituent quark models, without explicit pion cloud effects, the predictions for the magnetic moment, which dominates the transition, deviate necessarily from the data for low Q^2, showing that an effective pion cloud within the VMD current is not sufficient. Because of the opening of the πN channel, an explicit pion cloud contribution must be considered [4, 5]. Therefore our results in the physical low Q^2 region give only the constituent quark core contributions [4], usually labeled as "bare" contributions. Our results for these contributions are consistent with independent calculations of dynamical reaction models [19, 20] based on hadronic degrees of freedom. As for the E2 and C2 multipoles, in the large N_c limit [21] they represent second order corrections and are not considered here.

As pion cloud effects in lattice QCD are expected to be small for $m_\pi > 0.40$ GeV [22], the description of the lattice QCD data appears as the ideal testing ground for our constituent quark model, where pion cloud ef-

fects are not explicitly included. This defines the main goal of this work. At present, lattice calculations are still restricted to high pion masses $m_\pi > 350$ MeV, although technical improvements are increasingly allowing to reach nearer and nearer the physical region [23, 24]. In this work we took the quenched lattice QCD data from Refs. [24, 25]. In the unquenched calculations the effects of the sea quarks are explicitly considered. Although unquenched lattice QCD data is already available, it is not expected that the quenched and unquenched data differ substantially in the region $m_\pi > 0.40$ GeV [22]. In addition, there are some differences between unquenched data for the $\gamma N \to \Delta$ transition based on two different unquenched methods at similar pion masses [26] that still have to be clarified in the future. For these reasons we took the conservative option of using quenched data exclusively.

Our new model presented here is based on three specific features: i) the electromagnetic interaction with the constituent quarks is described within the impulse approximation, and considering VMD; ii) the wavefunctions of the quark-diquark system are parametrized by simple monopole factors reduced to the Hulthen form, with one or two effective range parameters that balance the details of the short range and the long range behavior of the system; iii) the constituent quark magnetic anomalous moment scales with the inverse of its mass, which we write, following [27] as a function of the current quarks mass.

A current based on VMD is suitable to describe the interaction of the photon with the constituent quark, for the quark-antiquark spin-1 vertex. Depending on the isospin, the intermediate meson pole corresponds to the ρ or the ω meson, at low Q^2, or, in the large Q^2 regime, to some other effective heavy meson pole M_h. The photon interaction is then described as proceeding through the production of an intermediate meson state which annihilates subsequently into a quark-antiquark pair. VMD is successful in the description of the electromagnetic interaction with nucleons.

II. EXTENSION OF THE QUARK MODEL TO THE LATTICE DATA REGION

The nucleon S-state wavefunction includes the correct spin structure for the quark-diquark spin 0 and 1 components, associated to isospin 0 and 1 states, respectively. For the Δ, since total isospin is 3/2, the S-state wavefunction reduces to the diquark spin 1 with isospin 1 structure. As for the scalar wavefunction, it takes the phenomenological form

$$\psi_B(P,k) = \frac{N_B}{m_s(\alpha_1+\chi_B)(\alpha_2+\chi_B)^{n_B}}, \qquad (1)$$

where $B = N, \Delta$, $\chi_B = \frac{(M-m_s)^2-(P-k)^2}{Mm_s}$, with M the baryon mass (for the nucleon or Δ), m_s is the diquark mass and N_B is a normalization constant. [In Ref. [4, 14] we considered $n_B = 1$ for the nucleon and $n_B = 2$ for the Δ]. The parameters α_i can be interpreted as Yukawa mass or range coefficients that distinguish between two different regimes for the momentum range.

The parametrization of the momentum dependence in terms of χ_B absorbs the dependence on the baryon masses M. The range parameters α_i for the nucleon and the Δ were fixed by the nucleon and $\gamma N \to \Delta$ form factor experimental data [4, 14]. In the spectator quark model the transition amplitude between a initial (momentum P_-) nucleon (N) and a final (momentum P_+) baryon $B = N, \Delta$ can be written, in impulse approximation [4, 5, 14] as

$$J^\mu = 3\sum_\lambda \int_k \bar{\Psi}_B(P_+,k) j_I^\mu \Psi_N(P_-,k), \qquad (2)$$

where $\Psi_B(P,k)$ is the complete baryon wavefunction (including spin, isospin and momentum dependence), λ is the diquark polarization and j_I^μ the quark current operator dependent of the hadron isospin I. The baryon polarizations are suppressed for simplicity. The symbol $\int_k$ represents the invariant integral in the on-shell diquark moment $\int_k \equiv \int \frac{d^3k}{2E_s(2\pi)^3}$, with E_s as the diquark energy. The factor 3 takes account for the isospin symmetrization.

The quark current [4, 5, 14] takes the general form

$$\begin{aligned} j_I^\mu \;=\; & \left(\frac{1}{6}f_{1+}(Q^2) + \frac{1}{2}f_{1-}(Q^2)\tau_3\right)\gamma^\mu + \\ & \left(\frac{1}{6}f_{2+}(Q^2) + \frac{1}{2}f_{2-}(Q^2)\tau_3\right)\frac{i\sigma^{\mu\nu}q_\nu}{2M_N}. \end{aligned} \qquad (3)$$

The constituent quark form factors $f_{i\pm}$ $(i = 1, 2)$ are parametrized using a VMD structure, and are normalized according to $f_{1\pm}(0) = 1$ and $f_{2\pm}(0) = \kappa_\pm$, where $\kappa_\pm$ are defined in terms of the quark anomalous magnetic moments κ_u and κ_d according with $\kappa_+ = 2\kappa_u - \kappa_d$ and $\kappa_- = \frac{1}{3}(2\kappa_u + \kappa_d)$. See Refs. [4, 14] for details.

In this first calculation the wavefunctions are reduced to S-states for the quark-diquark system. Although it is well known that angular momentum components besides the S-states are essential for the description of the nucleon in Deep Inelastic Scattering, their effects are not so evident in the elastic form factors, particularly for low Q^2 [14]. As for the Δ, calculations of the valence quark contribution associated with higher angular momentum states (D-states) suggest a small effect [5, 26]. In fact, even with only S-wave components in the wavefunctions for both the nucleon and the Δ, the model generates the dominant contribution for the $\gamma N \to \Delta$ transition, the dipole magnetic moment form factor. The non-zero contributions for the other form factors appear only when D states are included, or when pion cloud effects are taken in consideration [5, 26]. At the end, we will estimate the effect in the $\gamma N \to \Delta$ quadrupole transition form factors, in the lattice regime, from adding the D-states to the present model.

As shown in [14] the diquark mass scales out from the formulas obtained for the electromagnetic form factors. This allowed us to ignore any explicit dependence on the quark mass. This mass dependence was present in the nucleon, Δ and vector meson masses in an implicit way only. However, in order to grant a comparison of our results with the lattice data, we extend here the model to include a dependence of the quark anomalous moment κ_u and κ_d on the quark mass. This dependence was not explicitly considered in Refs. [4, 14] because they dealt only with the physical data.

Inspired by Ref. [27], which combines chiral symmetry with conventional quark models, and applies an analytic continuation of the chiral expansion to the simple SU(6) model, we then use the smooth variation of the hadronic properties with the current quark masses above 60 MeV. Therefore, in the spirit of the constituent quark models, we assume here that κ_q $(q = u, d)$ scales with $1/M_q$, where M_q is the constituent quark mass. Labeling the quark anomalous moment at the physical point by κ_q^0, we can write κ_q, for an arbitrary constituent quark mass M_q, as

$$\kappa_q = \frac{M_q^{\rm phy}}{M_q}\kappa_q^0, \tag{4}$$

where $M_q^{\rm phy}$ is the constituent quark mass at the physical point.

To include an explicit dependence on the quark mass we consider the parametrization due to Cloet *et. al.* [27]

$$M_q \;=\; M_\chi + c m_q \;=\; M_\chi + c m_q^{\rm phy}\frac{m_\pi^2}{(m_\pi^{\rm phy})^2}, \tag{5}$$

where m_q ($m_q^{\rm phy}$) is the (physical) current quark mass, M_χ is a new parameter corresponding to the constituent quark mass in the chiral limit ($m_q = 0$), and c is a coefficient of the order of the unity. In the same equation m_π stands for the pion mass in the model, a parameter to be fixed by the lattice data, and $m_\pi^{\rm phy}$ for the physical mass. Following Ref. [27] we considered $c m_q^{\rm phy} = 5.9$ MeV.

The parametrization in (5) is most sensitive to M_χ. Reference [27] fixes $M_\chi = 0.42$ GeV. Different descriptions using quark models and different lattice sizes lead to different values [15]. For this reason we use M_χ as the only free parameter allowed to vary in the calculation presented here. It is needed to introduce an explicit dependence of the nucleon magnetic moment on the pion mass, and to enable the connection of the constituent quark model to the lattice QCD calculations [13, 23, 24, 25, 28, 29].

We consider three different cases. First we considered the case $M_\chi = +\infty$, corresponding to the limit where κ_q has no dependence on m_π. We tested also the original parametrization $M_\chi = 0.42$ GeV [27]. Finally, although we did not performed a systematic fit, we tested several other values of M_χ.

Furthermore, to extend the spectator model to the region of the quenched lattice QCD data we still need to consider the nucleon and Δ masses determined by the lattice simulations [24]. This brings to the calculation an implicit dependence on the pion mass. As for the vectorial mesons included in the VMD, quark current picture, we use the parametrization [28]

$$m_\rho = c_0 + c_1 m_\pi^2, \tag{6}$$

where $c_0 = 0.776$ GeV and $c_1 = 0.427$ GeV^{-1}. The simple parametrization (6) describes well the quenched lattice QCD data and is consistent with finite volume corrections [13, 30]. In this applications we consider this parametrization together with model II of the Refs. [4, 14], where the heavy vectorial meson mass is $M_h = 2M_N$, with M_N the nucleon mass of the lattice calculations.

III. RESULTS

In Fig. 1 we compare the predictions of our model for the nucleon form factor with the lattice data from Ref. [23], corresponding to the three values of M_χ. We consider in particular the isovector nucleon form factor because the contributions of the disconnected diagrams vanishes in lattice calculations, if the flavor SU(2) symmetry is assumed as in Ref. [23]. We considered the lowest pion masses and the smallest lattice spacings. In particular, we take the data corresponding to $m_\pi = 504$ MeV, 649 MeV ($a = 0.059$ fm) and $m_\pi = 615$ MeV ($a = 0.078$ fm). For larger lattice spacings a dependence on a is observed [23]. From the figure we conclude that $M_\chi = 0.42$ GeV and 0.80 GeV gives an excellent description of the nucleon data.

Still, when the pion mass increases there is a systematic deviation of our model from the lattice data of Ref. [23]. This deviation may be a consequence of the fact that the wavefunction parametrization in terms of a low momentum scale (α_1^N) (long range behavior) and high momentum scale (α_2^N) was kept unchanged. As the pion masses vary, at least the long range parameter may vary, and may vary more rapidly for larger pion masses. As for the short range parameter (α_2^N) it sets the scale below which the short range physics becomes important. As observed in applications of the finite-range regularization effective field theory, the scale associated to the momentum cut-off (and consequently short range effects) can be expressed by an universal regulator which describes simultaneously the large pion mass regime and the physical regime [13, 29, 31]. In our calculation α_2^N is related to that universal regulator, and it is expected not to depend crucially on the pion mass value.

For a finer analysis of our results, we compare our predictions for the isovector nucleon magnetic moment in physical nucleon magnetons, with the lattice QCD data and the chiral result from Ref. [32]. To convert $G_M^v(0)$ given by the lattice data in units $\frac{e}{2M_N}$ (M_N nucleon

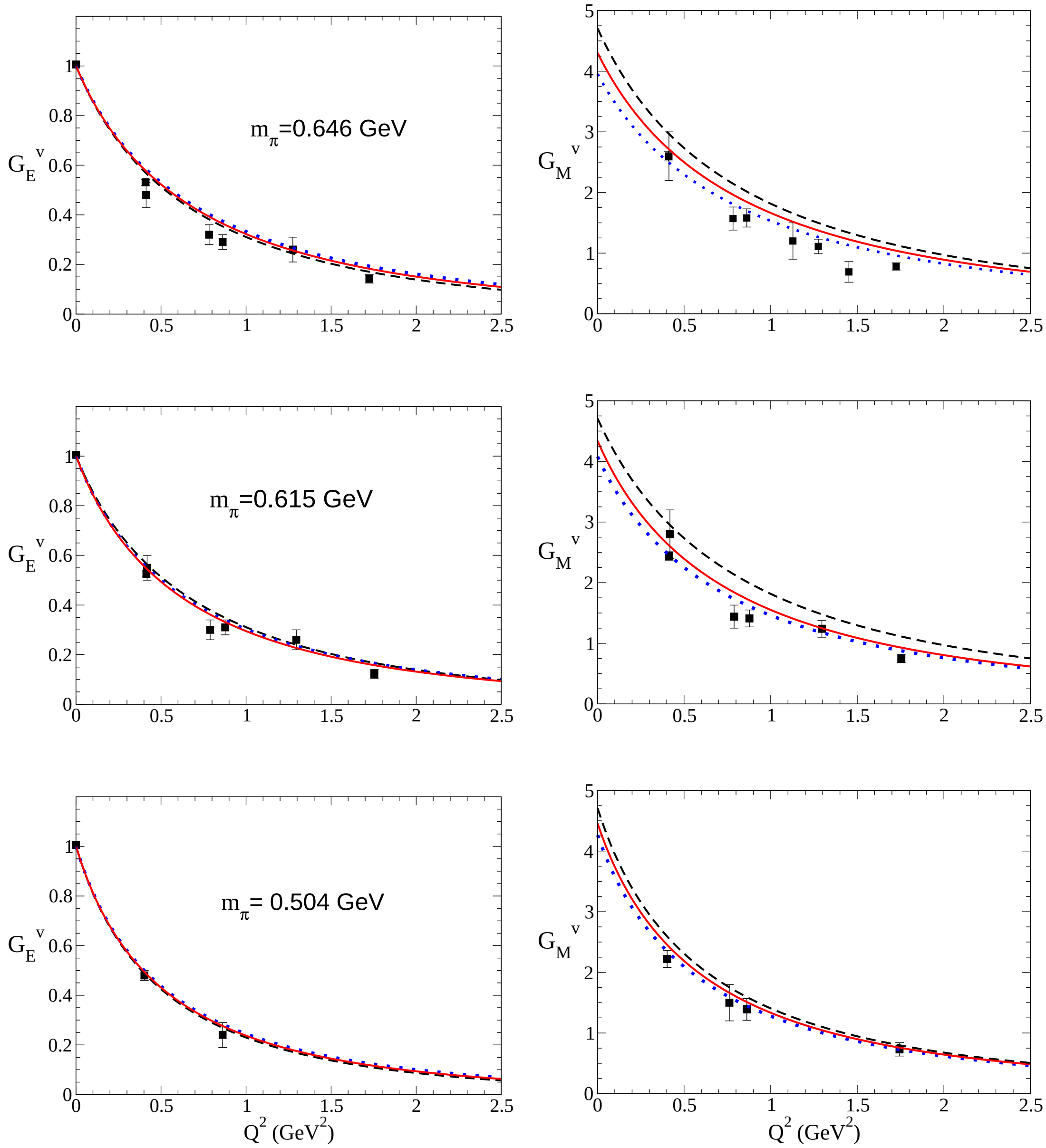

FIG. 1: Nucleon isovector form factors in lattice QCD from Ref. [23] and our predictions corresponding respectively to M_χ $+\infty$ (dashed line), 0.42 GeV (dotted line) and 0.80 GeV (solid line).

mass in lattice) to "physical" units $\frac{e}{2M_N^{\rm phy}}$, we need to use the transformation $G_M^v(0) \to \frac{M_N^{\rm phy}}{M_N} G_M^v(0)$. The results are presented in the Fig. 2 for $M_\chi = 0.42$ GeV and $M_\chi = 0.80$ GeV, as function of m_π. The agreement of our results with the chiral expression from Ref. [32] shows the consistency of our calculations with chiral calculations.

Finally, in Fig. 3 we compare the quenched lattice data for the $\gamma N \to \Delta$ transition from Ref. [24] with the spectator quark model corresponding to $M_\chi = 0.42$ GeV, 0.80 GeV and $M_\chi = +\infty$. All cases shown correspond to a lattice spacing of $a = 0.092$ fm. The contribution of the quark core extracted from [20] at the physical point is also included. To be consistent, we considered the parametrization of Eq. (6) for m_ρ, although the original Ref. [24] gives slightly different results. As for M_N and M_Δ, we use the values derived directly from the lattice data [24]. For the larger pion masses we observe an almost perfect agreement between the predictions of $M_\chi = +\infty$ and the data, although $M_\chi = 0.80$ GeV is also close. For $m_\pi = 0.411$ GeV we have also a good agreement, except for a slight deviation from the lattice data for 0.7 GeV$^2 < Q^2 < 2$ GeV^2. This deviation may result from the effect of the pion cloud for light pions, predicted to be important for $m_\pi < 0.40$ GeV [22]. As for the physical pion mass case, our model is coherent

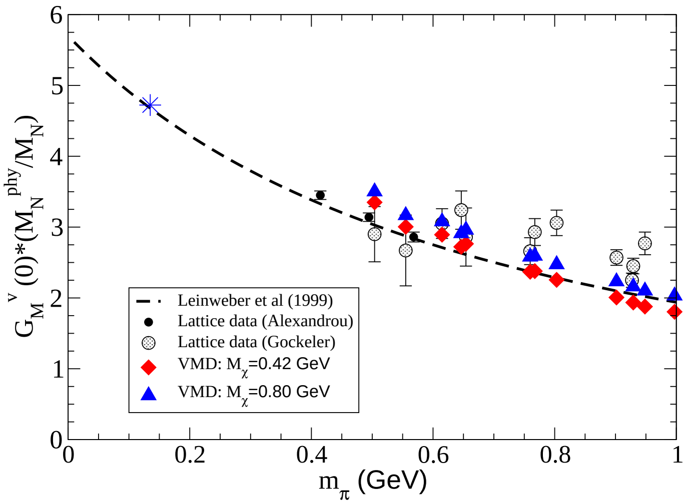

FIG. 2: Nucleon isovector magnetic moment as function of m_π, in physical nucleon magnetons. Theoretical line from Ref. [32]. Our results are labeled "VMD" and shown for two values of the M_χ parameter.

with the constituent quark core data labeled 'Bare Form Factor', which is extracted indirectly from experiment [4, 20]. In this case all lines coincide.

As mentioned already, the D-states in the Δ wavefunction induce contributions for the subleading electric and Coulomb quadrupole transition form factors G_E^* and G_C^*. The exact contributions depend on the specific parametrization, in particular on the admixture coefficients for the two D-waves. In Ref. [26] it was shown how a percentage of D-states of $\approx 0.9\%$ can provide an excellent description of the quadrupole lattice data, without significantly changing the description of the dominant magnetic dipole form factor. That application corresponds to the $M_\chi = +\infty$ limit and estimates the effect of the valence quark (or bare) contributions as only $\approx$ 20% of the total quadrupole for both G_E^* and G_C^*, at the physical point (the remaining contribution being the pion cloud). On the other hand, in that particular parametrization, and in the region of study $Q^2 < 1.5$ GeV^2, the effect of the constituent quark mass dependence described by Eq. (4) corresponds to a correction of less than 20% for both form factors, a correction smaller than the typical lattice errorbands ($\approx$ 30%).

IV. CONCLUSIONS

Chiral perturbation methods for lattice QCD extrapolations are useful for small Q^2 (like $Q^2 < 1.5$ GeV^2), but are not adequate for the high Q^2 region. Constituent quark models can supply an alternative guidance for lattice QCD extrapolations. In this work we consider a covariant constituent quark model of the nucleon and the Δ based on the spectator formalism for the quark-diquark system (covariance is an important issue for the description of the high momentum transfer processes).

The model contains no explicit pion cloud, besides the effects included in the ρ term of the VDM picture for the electromagnetic current, and is therefore reduced to the bare quark hadron structure. Since pion cloud effects are expected to be negligible for large values of the pion mass, the comparison of the model to the quenched data for pion masses larger than 450 MeV is justified a priori. To accomplish that comparison the initial covariant quark model was extended to the lattice data region, $m_\pi < 700$ MeV. In light of the work in Ref. [27], this extension was done by introducing the constituent quark mass in the chiral limit parameter, M_χ, and the nucleon ρ, Δ masses used in the lattice calculations.

The main conclusion is that the covariant constituent quark model which was previously calibrated by means of a quantitative description of the nucleon, $\gamma N \to \Delta$ and Δ form factor data in the physical region, as shown in previous works [4, 14, 17], after a simple extension involving one parameter only, describes also quantitatively well the lattice results for the nucleon isovector form factor and the $\gamma N \to \Delta$ "bare" magnetic form factor. With $M_\chi = 0.42$ GeV, consistently with the range $M_\chi = 0.3 - 0.45$ GeV suggested by several constituent quark models [6, 7, 9, 11], we obtain a very good description of the lattice nucleon form factors data, but underestimate the $\gamma N \to \Delta$ magnetic moment form factor data by less than 9% at low Q^2 for $m_\pi < 500$ MeV. Note that the original parametrization $M_\chi = 0.42$ GeV [27] was a result of a phenomenological fit and was not derived from first principles.

An optimal description of the lattice data for both nucleon and $\gamma N \to \Delta$ transition form factors is achieved, once the scale of the constituent quark mass in the chiral limit is fixed as $M_\chi = 0.80$ GeV. [The values $M_\chi = 0.42$ and 0.80 GeV are better for the nucleon data; $M_\chi = 0.80$ GeV and $+\infty$ are better for the $\gamma N \to \Delta$ data]. The parameter $M_\chi = 0.80$ GeV is relatively large when compared with alternative constituent quark models, which may indicate that some of the effective quark-antiquark configurations contained in our model through the VMD mechanism in the electromagnetic current, for instance, do not correspond to the lattice calculations in the quenched approximation. Still, the possibility of adjusting other parameters is very promising.

A refinement of the description can be pursued, by increasing the number of adjustable parameters, or simply by introducing an extra dependence on the pion mass. For instance, in the future we certainly plan to check the assumption that the α_1 parameter which controls the long range regime does not depend on the pion mass value. In addition, we may consider as well an explicit dependence of the heavier meson mass pole M_h, that regulates the short range effect in VMD mechanism, on the pion mass. Alternatively, a pion mass dependence of the α_2 parameter can be introduced. Furthermore, we want to study the quality of the description in the high pion

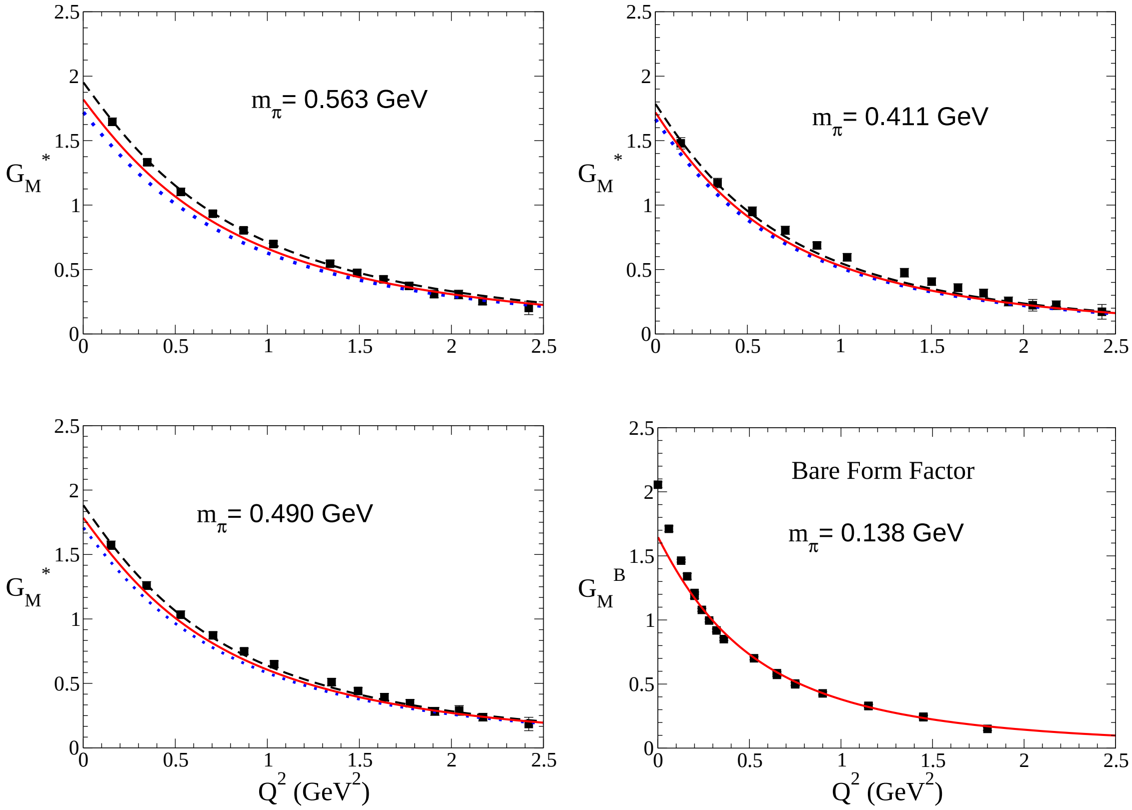

FIG. 3: $\gamma N \to \Delta$ magnetic form factor lattice data from Ref. [24] and 'bare' form factors from Ref. [20] (physical point) compared with models $M_\chi = 0.42$ GeV, 0.80 GeV and $M_\chi = +\infty$. The conventions are the same as in Fig. 1.

masses region, not probed yet here.

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[1] V. Pascalutsa and M. Vanderhaeghen, Phys. Rev. D **73**, 034003 (2006) [arXiv:hep-ph/0512244].

[2] T. A. Gail and T. R. Hemmert, Eur. Phys. J. A **28**, 91 (2006) [arXiv:nucl-th/0512082].

[3] V. Pascalutsa, M. Vanderhaeghen and S. N. Yang, Phys. Rept. **437**, 125 (2007) [arXiv:hep-ph/0609004].

[4] G. Ramalho, M. T. Peña and F. Gross, Eur. Phys. J. A 36, 329 (2008) arXiv:0803.3034 [hep-ph].

[5] G. Ramalho, M. T. Peña and F. Gross, Phys. Rev. D **78**, 114017 (2008) [arXiv:0810.4126 [hep-ph]].

[6] F. Coester and D. O. Riska, Nucl. Phys. A **728**, 439 (2003) [arXiv:nucl-th/0306002]; F. Coester, K. Dannbom and D. O. Riska, Nucl. Phys. A **634**, 335 (1998) [arXiv:hep-ph/9711458].

[7] B. Julia-Diaz, D. O. Riska and F. Coester, Phys. Rev. C **69**, 035212 (2004) [Erratum-ibid. C **75**, 069902 (2007)] [arXiv:hep-ph/0312169].

[8] S. Boffi, L. Y. Glozman, W. Klink, W. Plessas, M. Radici and R. F. Wagenbrunn, Eur. Phys. J. A **14**, 17 (2002) [arXiv:hep-ph/0108271]; R. F. Wagenbrunn, S. Boffi, W. Klink, W. Plessas and M. Radici, Phys. Lett. B **511**, 33 (2001) [arXiv:nucl-th/0010048].

[9] M. M. Giannini, E. Santopinto and A. Vassallo, Eur. Phys. J. A **12**, 447 (2001) [arXiv:nucl-th/0111073].

[10] E. Pace, G. Salme, F. Cardarelli and S. Simula, Nucl. Phys. A **666**, 33 (2000) [arXiv:nucl-th/9909025]; F. Cardarelli and S. Simula, Phys. Rev. C **62**, 065201 (2000) [arXiv:nucl-th/0006023].

[11] C. V. Christov *et al.*, Prog. Part. Nucl. Phys. **37**, 91 (1996) [arXiv:hep-ph/9604441].

[12] D. Merten, U. Loring, K. Kretzschmar, B. Metsch and H. R. Petry, Eur. Phys. J. A **14**, 477 (2002) [arXiv:hep-ph/0204024].

[13] R. D. Young, D. B. Leinweber and A. W. Thomas, Phys. Rev. D **71**, 014001 (2005) [arXiv:hep-lat/0406001].

[14] F. Gross, G. Ramalho and M. T. Peña, Phys. Rev. C **77**, 015202 (2008) [arXiv:nucl-th/0606029]; F. Gross, G. Ramalho and M. T. Peña, Phys. Rev. C **77**, 035203 (2008).

[15] H. H. Matevosyan, G. A. Miller and A. W. Thomas, Phys. Rev. C **71**, 055204 (2005).

[16] H. H. Matevosyan, A. W. Thomas and G. A. Miller, Phys. Rev. C **72**, 065204 (2005) [arXiv:nucl-th/0508049].

[17] G. Ramalho and M. T. Pena, J. Phys. G **36**, 085004 (2009) [arXiv:0807.2922 [hep-ph]]; G. Ramalho, M. T. Pena and F. Gross, Phys. Lett. B **678**, 355 (2009) [arXiv:0902.4212 [hep-ph]].

[18] U. G. Meissner, AIP Conf. Proc. **904**, 142 (2007) [arXiv:nucl-th/0701094]; N. Kaiser, Phys. Rev. C **68**, 025202 (2003) [arXiv:nucl-th/0302072].

[19] T. Sato and T. S. H. Lee, Phys. Rev. C **63**, 055201 (2001) [arXiv:nucl-th/0010025]; S. S. Kamalov, S. N. Yang, D. Drechsel, O. Hanstein and L. Tiator, Phys. Rev. C **64**, 032201(R) (2001) [arXiv:nucl-th/0006068].

[20] B. Julia-Diaz, T. S. H. Lee, T. Sato and L. C. Smith, Phys. Rev. C **75**, 015205 (2007) [arXiv:nucl-th/0611033].

[21] E. E. Jenkins, X. Ji and A. V. Manohar, Phys. Rev. Lett. **89**, 242001 (2002) [arXiv:hep-ph/0207092]; V. Pascalutsa and M. Vanderhaeghen, Phys. Rev. D **76**, 111501 (2007) [arXiv:0711.0147 [hep-ph]].

[22] J. D. Ashley, D. B. Leinweber, A. W. Thomas and R. D. Young, Eur. Phys. J. A **19**, 9 (2004) [arXiv:hep-lat/0308024].

[23] M. Gockeler, T. R. Hemmert, R. Horsley, D. Pleiter, P. E. L. Rakow, A. Schafer and G. Schierholz [QCDSF Collaboration], Phys. Rev. D **71**, 034508 (2005) [arXiv:hep-lat/0303019].

[24] C. Alexandrou, G. Koutsou, H. Neff, J. W. Negele, W. Schroers and A. Tsapalis, Phys. Rev. D **77**, 085012 (2008) [arXiv:0710.4621 [hep-lat]].

[25] C. Alexandrou, G. Koutsou, J. W. Negele and A. Tsapalis, Phys. Rev. D **74**, 034508 (2006) [arXiv:hep-lat/0605017].

[26] G. Ramalho and M. T. Pena, Phys. Rev. D **80**, 013008 (2009) [arXiv:0901.4310 [hep-ph]].

[27] I. C. Cloet, D. B. Leinweber and A. W. Thomas, Phys. Rev. C **65**, 062201 (2002) [arXiv:hep-ph/0203023].

[28] D. B. Leinweber, A. W. Thomas, K. Tsushima and S. V. Wright, Phys. Rev. D **64**, 094502 (2001) [arXiv:hep-lat/0104013].

[29] P. Wang, D. B. Leinweber, A. W. Thomas and R. D. Young, Phys. Rev. D **75**, 073012 (2007) [arXiv:hep-ph/0701082].

[30] C. R. Allton, W. Armour, D. B. Leinweber, A. W. Thomas and R. D. Young, PoS **LAT2005**, 049 (2006) [arXiv:hep-lat/0511004].

[31] A. W. Thomas, P. A. M. Guichon, D. B. Leinweber and R. D. Young, Prog. Theor. Phys. Suppl. **156**, 124 (2004) [arXiv:nucl-th/0411014].

[32] D. B. Leinweber, D. H. Lu and A. W. Thomas, Phys. Rev. D **60**, 034014 (1999) [arXiv:hep-lat/9810005].